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Gluonium as Bound State of Massive Gluons Described by the Joos-Weinberg Wave Functions.

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ABSTRACT. On the basis of the Kadyshevsky equal-time (quasipotential) approach, a set of partial-wave equations is derived for the wave function of a gluonium, a bound state of two gluons. The field operators of constituent gluons are considered as six component quantities according to the Joos-Weinberg $2(2S+1)$- component approach. The quasiclassical quantization condition for relativistic two-particle states and the above set can be used for calculations of gluonium energy levels.

RESUMEN. Partiendo del método de tiempos iguales (cuasi-potencial) de Kadyshevsky se deduce un conjunto de ecuaciones de ondas parciales para la función de onda del gluonio, el estado ligado de dos gluones. Los operadores de campo de los gluones constituyentes se consideran cantidades de seis componentes de acuerdo con el modelo de Joos-Weinberg de 2 (2S + 1) componentes. La condición de cuantización cuasi-clásica para los estados relativistas de dos partículas y el conjunto de ecuaciones mencionado pueden emplearse para calcular los niveles de energía del gluonio.

KEYWORDS: quantum chromodynamics, equal-time (quasipotential) approach, gluonium (glueball), Joos-Weinberg approach
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1 Introduction

The existence of a gluonium, which is a color singlet of the bound state of two or more gluons, is predicted by all the models of quantum chromodynamics (QCD), the lattice models [1], the sum rules [2], the bag model [3] and the effective Lagrangian approach [4]. Experimental searches of these states are in progress, see, e. g., for the reviews in ref. [5]. It follows from the analysis of obtained results that the most likely candidates for glueballs are the following meson resonances [6]:

Table. *The eventual candidates for glueball.*

	References	J^{PC}	$n\,^{2S+1}L_J$	Reactions of observations
$\sigma(750)$	[7]	0^{++}	$1\,^1S_0$	$\pi \pm N_{pol} \to \sigma N$
$\imath(1460)$	[8]	0^{-+}	$1\,^3P_0$	$J/\psi \to \gamma X$
$G(1590)$	[9]	0^{++}	$2\,^1S_0$	$\pi^- p \to \eta\eta n$
$\vartheta(1720)$	[10, 12]	2^{++}	$2\,^5S_2$	$J/\psi \to \gamma X$
$g(2050)$	[11]-[14]	2^{++}	$3\,^5S_2$	$\pi^- p \to \phi\phi n, \quad J/\psi \to \gamma X$
$\xi(2220)$	[11]	2^{++}	$2\,^3D_2 \quad or \quad 2\,^5D_2$	$J/\psi \to \gamma X$
$g'(2300)$	[11]-[14]	2^{++}	$3\,^5D_2$	$\pi^- p \to \phi\phi n, \quad J/\psi \to \gamma X$
$g''(2350)$	[11]-[14]	2^{++}	$3\,^1D_2$	$\pi^- p \to \phi\phi n, \quad J/\psi \to \gamma X$

In connection with that it is very important to describe the gluonium spectra theoretically. Attempts have been made earlier to consider gluonium in the framework of the potential model with massive structure gluons [6, 15] analogous to the non-relativistic description of the quark-antiquark system, ref. [16].

At present, the relativistic single-time approach [17]-[19] is used widely for the description of two-particle systems like quarkonium. The necessity of allowance for relativistic effects is caused by the fact that in many cases the contribution of the relativistic corrections is of the same order as the contribution of the non-relativistic Hamiltonian. In the present work, the quasipotential equal-time approach is employed for the description of two-gluon bound states [1] consisting of the structural gluons which are described by the six-component Weinberg's wave functions [21]-[24]. The results of the papers [25, 26], devoted to the covariant three-dimensional description of the composite system formed by two particles with the $S = 1$ spin, are used. The remarkable feature of our formalism is the locality of the corresponding quasipotential in the Lobachevsky momentum space. This is achieved by the separation of the kinematical Wigner rotations and "resetting" all spin indices to the one momentum, for details see refs. [18, 27, 28]. Moreover, the quasipotential for interaction of two vector particles is the same as the quasipotential for interaction of two spinor particles with corresponding substitutions accounting for the spin difference and the normalization.

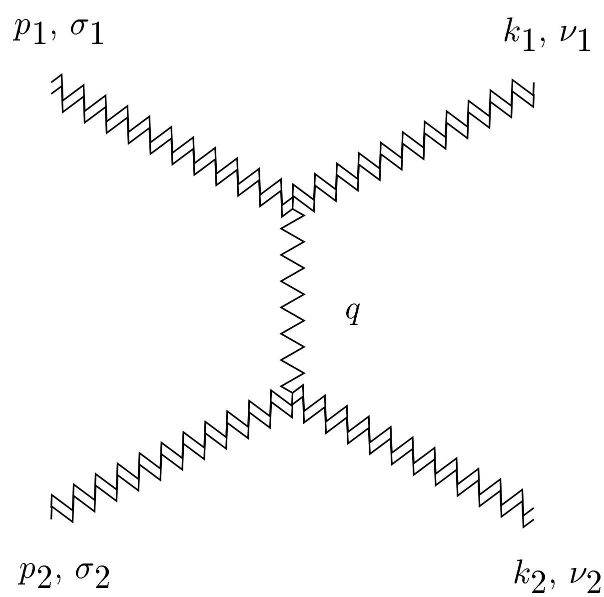
[1] The two-gluon bound system has recently been described on the basis of the Bethe-Salpeter equation, ref. [20].

The quasipotential equation with the one-boson exchange potential, obtained in Section II, is reduced to the finite-difference partial-wave equations in the relativistic configurational representation (RCR) in Section III. As was shown earlier [29], it is possible to develop the relativistic analog of the WKB methods in the RCR. This method has been successfully used to find the quarkonium mass spectrum [19]. The use of the quasiclassical quantization condition for relativistic two-particle states [19, a], see Section IV, also allows calculation the gluonium energy levels.

2 Spin structure of the relativistic potential for two-gluon interaction in the momentum representation

At present, the gluon could be described as a massive particle with dynamical mass appearing due to the existence of color charge and self-interaction. This fact permits one to eliminate some contradictions in the results of calculations of the proton formfactor and the effective coupling constant $\alpha_S(q^2)$ on the basis of QCD (see in this connection ref. [30]).

Therefore, we begin by considering the quasipotential for two-gluon interaction in the momentum representation as that of gluonium consisting of the structure massive gluons with the intermediate interaction of the gauge massless gluon. The corresponding diagram describing this process is drawn at *Fig. 1* . The Feynman matrix element $< p_1, p_2; \sigma_1, \sigma_2 \mid \hat{T}^{(2)} \mid k_1, k_2; \nu_1, \nu_2 >$ corresponding to this diagram is considered to be the quasipotential, $\hat{V}^{(2)} = \hat{T}^{(2)}$.



To find the form of the quasipotential it is necessary to know the Feynman rules for the vertices of interaction of the structure gluons and the massless gauge gluon.

In ref. [21] the attractive $2(2S+1)$- formalism for the description of particles of spin $S = 1$ has been proposed. As opposed to the Proca functions, which transform according

to the $(\frac{1}{2},\frac{1}{2})$ representation of the Lorentz group in the case of $S=1$, the spinor functions are constructed via the representation $(S,0)\oplus(0,S)$ in the Joos-Weinberg formalism. This way of description of higher spin particles is on an equal footing to the Dirac description of spinor particles whose wave functions transform according to the $(\frac{1}{2},0)\oplus(0,\frac{1}{2})$ representation. The $2(2S+1)$- component analogues of the Dirac functions in the momentum space are[2]

$$U(\vec{p}) = \frac{1}{\sqrt{2}}\begin{pmatrix} D^S\left(\alpha(\vec{p})\right)\xi_\sigma \\ D^S\left(\alpha^{-1\,+}(\vec{p})\right)\xi_\sigma \end{pmatrix}, \tag{2.1}$$

for the positive-energy states; and

$$V(\vec{p}) = \frac{1}{\sqrt{2}}\begin{pmatrix} D^S\left(\alpha(\vec{p})C^{-1}\right)\xi^*_\sigma \\ D^S\left(\alpha^{-1\,+}(\vec{p})C^{-1}\right)(-1)^{2S}\xi^*_\sigma \end{pmatrix}, \tag{2.2}$$

for the negative-energy states, with the following notations:

$$\alpha(\vec{p}) = \frac{p_0 + M + (\vec{\sigma}\vec{p})}{\sqrt{2M(p_0+M)}}, \quad C = -i\sigma_2; \tag{2.3}$$

and $D^S(A) \equiv D^{(S,0)}(A)$ is the Lorentz group representation by matrices with $(2S+1)$ rows and columns[3]. In the case of $S=1$, one has

$$D^{(1,0)}\left(\alpha(\vec{p})\right) = 1 + \frac{(\vec{S}\vec{p})}{M} + \frac{(\vec{S}\vec{p})^2}{M(p_0+M)}. \tag{2.4}$$

In spite of some antiquity of this formalism, in our opinion, it does not deserve to be retired. Recently, some attention has been paid to this formalism [22]-[26].

In the articles [21, b,h-j] the Feynman diagram technique is discussed for the vector particles in the above-mentioned six-component formalism for quantum electrodynamics (QED). The following Lagrangian:

$$\begin{aligned}\mathcal{L}^{QED} &= \bar{\Psi}(x)\Gamma_{\mu\nu}\bar{\nabla}_\mu\vec{\nabla}_\nu\Psi(x) - M^2\bar{\Psi}(x)\Psi(x) - \frac{1}{4}F_{\mu\nu}F_{\mu\nu} + \\ &+ \frac{e\lambda}{12}F_{\mu\nu}\bar{\Psi}(x)\gamma_{5,\mu\nu}\Psi(x) + \frac{e\kappa}{12M^2}\partial_\alpha F_{\mu\nu}\bar{\Psi}(x)\gamma_{6,\mu\nu,\alpha\beta}\nabla_\beta\Psi(x)\end{aligned} \tag{2.5}$$

has been used there. In the above formula we have $\nabla_\mu = -i\partial_\mu - eA_\mu$; $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the electromagnetic field tensor; A_μ is the 4- vector of the electromagnetic field; $\bar{\Psi}, \Psi$ are the six-component wave functions (WF) of the massive vector particle. The following expression has been obtained for the interaction vertex of a vector particle with a photon [21, j],[26]:

$$-e\Gamma_{\alpha\beta}(p+k)_\beta - \frac{ie\lambda}{6}\gamma_{5,\alpha\beta}q_\beta + \frac{e\kappa}{6M^2}\gamma_{6,\alpha\beta,\mu\nu}q_\beta q_\mu(p+k)_\nu, \tag{2.6}$$

where $\Gamma_{\alpha\beta} = \gamma_{\alpha\beta} + \delta_{\alpha\beta}$; $\gamma_{\alpha\beta}$; $\gamma_{5,\alpha\beta}$; $\gamma_{6,\alpha\beta,\mu\nu}$ are $6\otimes 6$- matrices which have been considered in ref. [21, b,g]

[2]These functions obeys the orthonormalization equations, $U^+(\vec{p})\gamma_{44}U(\vec{p}) = 1$ and analogous equation exists for $V(\vec{p})$, the functions of negative-energy states.

[3]The technique of construction of $D^S(A)$ could be found in [21, k].

$$\gamma_{ij} = \begin{pmatrix} 0 & \delta_{ij} - S_i S_j - S_j S_i \\ \delta_{ij} - S_i S_j - S_j S_i & 0 \end{pmatrix}$$

$$\gamma_{i4} = \gamma_{4i} = \begin{pmatrix} 0 & iS_i \\ -iS_i & 0 \end{pmatrix}, \; \gamma_{44} = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix},$$

(here, S_i are the spin matrices for a vector particle),

$$\gamma_{5,\alpha\beta} = i[\gamma_{\alpha\mu}, \gamma_{\beta\mu}], \tag{2.7}$$

$$\gamma_{6,\alpha\beta,\mu\nu} = [\gamma_{\alpha\mu}, \gamma_{\beta\nu}]_+ + 2\delta_{\alpha\mu}\delta_{\beta\nu} - [\gamma_{\beta\mu}, \gamma_{\alpha\nu}]_+ - 2\delta_{\beta\mu}\delta_{\alpha\nu}, \tag{2.8}$$

e is electron charge, λ and κ are the quantities which correspond to the magnetic dipole moment and the electric quadrupole moment, respectively; M is the vector particle mass.

In the case of interaction of two structure gluons we also use the $2(2S+1)$- component formalism. Since the gluon $2(2S+1)$- dimensional WF can not be directly introduced into the Lagrangian by means of the standard procedure of lengthening the derivative (covariantization), we add the terms defining the structure gluons manually:

$$\mathcal{L}^{QCD}_{aux} = \bar{g}\bar{\nabla}_\mu \gamma^{\mu\nu} \vec{\nabla}_\nu g - M_g^2 \bar{g} g \tag{2.9}$$

into the commonly used QCD Lagrangian:

$$\mathcal{L}^{QCD} = i\bar{q}\gamma^\mu \nabla_\mu q - m_q \bar{q} q - \frac{1}{4} G^a_{\mu\nu} G^{a,\mu\nu}. \tag{2.10}$$

Where

$$\nabla_\mu g(x) = \partial_\mu g - if T^a_{adj} B^a_\mu g \tag{2.11}$$

and $\bar{g}$, g are the $2(2S+1)$- dimensional WF of the color octet, T^a_{adj} are the $SU(3)$ generators in the adjoin representation. Employing the technique of functional integration and using a scheme which is analogous to the $\bar{q}qg$- vertex case, we suggest that the interaction vertex of two structure gluons with the gauge gluon has the following analytical form[4]:

$$if\gamma^{\mu\nu}(p+k)_\nu (T^a_{adj})_{\alpha\beta} \tag{2.12}$$

without taking into account the multipole momenta (compare with the QED expression (2.6)).

Using results of refs. [25, 28], let us represent the Feynman matrix element corresponding to the diagram of one-gluon exchange (see *Fig. 1*) as:

$$\begin{aligned} &< p_1, p_2; \sigma_1, \sigma_2 | \hat{V}^{(2)} | k_1, k_2; \nu_1, \nu_2 > = < p_1, p_2; \sigma_1, \sigma_2 | \hat{T}^{(2)} | k_1 . k_2; \nu_1, \nu_2 > = \\ &= \sum_{\sigma_{ip}, \nu_{ip}, \nu_{ik} = -1}^{1} D^{+\,(S=1)}_{\sigma_1 \sigma_{1p}} \left\{ V^{-1}(\Lambda_{\mathcal{P}}, p_1) \right\} D^{+\,(S=1)}_{\sigma_2 \sigma_{2p}} \left\{ V^{-1}(\Lambda_{\mathcal{P}}, p_2) \right\} \times \\ &\times \; V^{\nu_{1p}\nu_{2p}}_{\sigma_{1p}\sigma_{2p}}(\vec{k}(-)\vec{p}, \vec{p}) D^{(S=1)}_{\nu_{1p}\nu_{1k}} \left\{ V^{-1}(\Lambda_{p_1}, k_1) \right\} D^{(S=1)}_{\nu_{1k}\nu_1} \left\{ V^{-1}(\Lambda_{\mathcal{P}}, k_1) \right\} \times \\ &\times \; D^{(S=1)}_{\nu_{2p}\nu_{2k}} \left\{ V^{-1}(\Lambda_{p_2}, k_2) \right\} D^{(S=1)}_{\nu_{2k}\nu_2} \left\{ V^{-1}(\Lambda_{\mathcal{P}}, k_2) \right\}, \end{aligned} \tag{2.13}$$

[4] We did not allow for the term $\delta_{\mu\nu}(p+k)_\nu$ which is the result of the auxiliary condition (Klein-Gordon equation) to the Weinberg equation. In the Kadyshevsky's approach all the particles, even in the intermediate states, are on the mass shell.

where

$$V^{\nu_{1p}\nu_{2p}}_{\sigma_{1p}\sigma_{2p}}(\vec{k}(-)\vec{p},\vec{p}) = \xi_{\sigma_{1p}}\xi_{\sigma_{2p}}\hat{V}^{(2)}(\vec{k}(-)\vec{p},\vec{p})\xi_{\nu_{1p}}\xi_{\nu_{2p}}. \tag{2.14}$$

After some calculation, using the formulae of ref. [25, 27]

$$U_\sigma(\vec{p}) = S_{\vec{p}}U_\sigma(\vec{0}), \quad S^{-1}_{\vec{p}}S_{\vec{k}} = S_{\vec{k}(-)\vec{p}} \cdot I \otimes D^1\left(V^{-1}(\Lambda_{\vec{p}},\vec{k})\right), \tag{2.15}$$

$$S^{-1}_{\vec{p}}\gamma_{\mu\nu}p_\nu S_{\vec{p}} = \gamma_{44}\left(p_\mu - \gamma_5 W_\mu(\vec{p})\right), \tag{2.16}$$

$$W_\mu(\vec{p}) \cdot D\left(V^{-1}(\Lambda_{\vec{p}},\vec{k})\right) = D\left(V^{-1}(\Lambda_{\vec{p}},\vec{k})\right) \cdot \left\{W_\mu(\vec{k}) + \frac{p_\mu + k_\mu}{M(\Delta_0 + M)}p_\nu W_\nu(\vec{k})\right\}, \tag{2.17}$$

$$k_\mu W_\mu(\vec{p}) \cdot D\left(V^{-1}(\Lambda_{\vec{p}},\vec{k})\right) = -D\left(V^{-1}(\Lambda_{\vec{p}},\vec{k})\right) \cdot p_\mu W_\mu(\vec{k}), \tag{2.18}$$

we come to the 4- current of vector particle[5]:

$$j^{\sigma_p\nu_p}_\mu(\vec{p},\vec{k}) = -f\xi_{\sigma_p}\left\{(p+k)_\mu + \frac{1}{M}W_\mu(\vec{p})(\vec{S}\vec{\Delta}) - \frac{1}{M}(\vec{S}\vec{\Delta})W_\mu(\vec{p})\right\}\xi_{\nu_p} \tag{2.19}$$

and to the quasipotential

$$\begin{aligned}\hat{V}^{(2)}(\vec{k}(-)\vec{p},\vec{p}) = -3f^2 &\left\{\frac{\left[p_0(\Delta_0+M) + (\vec{p}\vec{\Delta})\right]^2 - M^3(\Delta_0+M)}{M^3(\Delta_0 - M)} + \right.\\ + \quad & \frac{i(\vec{S}_1+\vec{S}_2)\left[\vec{p}\vec{\Delta}\right]}{\Delta_0 - M}\left[\frac{p_0(\Delta_0+M)+\vec{p}\vec{\Delta}}{M^3}\right] + \frac{(\vec{S}_1\vec{\Delta})(\vec{S}_2\vec{\Delta}) - (\vec{S}_1\vec{S}_2)\vec{\Delta}^2}{2M(\Delta_0 - M)} - \\ - \quad & \left.\frac{1}{M^3}\frac{\vec{S}_1\left[\vec{p}\vec{\Delta}\right]\vec{S}_2\left[\vec{p}\vec{\Delta}\right]}{\Delta_0 - M}\right\}.\end{aligned} \tag{2.20}$$

As used in the earlier works [18, 28], we have

$$\vec{\Delta} = \Lambda^{-1}_{\vec{p}}\vec{k} = \vec{k}(-)\vec{p} = \vec{k} - \frac{\vec{p}}{M}(k_0 - \frac{\vec{k}\vec{p}}{p_0 + M}), \tag{2.21}$$

$$\Delta_0 = (\Lambda^{-1}_{\vec{p}}k)_0 = (k_0p_0 - \vec{k}\vec{p})/M, \tag{2.22}$$

where $\vec{p}$, $\vec{k}$ are the covariant generalizations[6] of the vectors of particle momenta in c.m.s., before $\vec{p}_1 = -\vec{p}_2 = \vec{p}$ and after $\vec{k}_1 = -\vec{k}_2 = \vec{k}$ interaction; ξ^*, ξ are the analogues of Pauli spinors and $D^J_{\alpha\beta}$ is the Wigner matrix of the irreducible representation of the rotation group, which has dimension equal to $(2S+1)$ with the following form:

$$\begin{aligned}D^{(S=1)}\left\{V^{-1}(\Lambda_{\vec{p}},\vec{k})\right\} = &\frac{1}{2M(p_0+M)(k_0+M)(\Delta_0+M)}\left\{\left[\vec{p}\vec{k}\right]^2 + \right.\\ + \quad & \left[(p_0+M)(k_0+M) - \vec{k}\vec{p}\right]^2 - 2i\left[(p_0+M)(k_0+M) - \vec{k}\vec{p}\right]\vec{S}\left[\vec{p}\vec{k}\right] - \\ - \quad & \left. 2\{\vec{S}\left[\vec{p}\vec{k}\right]\}^2\right\}.\end{aligned} \tag{2.23}$$

[5] W_μ is the Pauli-Lyuban'sky 4- vector of relativistic spin.

[6] We omit the circles above the covariant generalizations of the momenta, as opposed to [25, 26, 28].

The expression (2.20) shows the advantages of the $2(2S+1)$- formalism, since it looks like the quasipotential for the interaction of two spinor particles with the substitutions $\frac{1}{2m(\Delta-m)} \Rightarrow \frac{1}{\vec{\Delta}^2}$ and $\vec{S} \Rightarrow \vec{\sigma}$.

3 System of two-particle partial-wave equations in the relativistic configurational representation

The transformations into the relativistic configurational representation (RCR) have the following form:

$$V(r,\vec{n};\vec{p}) = \frac{1}{(2\pi)^3}\int d\Omega_\Delta\, \xi^*(\vec{\Delta};\vec{r})V(\vec{\Delta},\vec{p}), \tag{3.1}$$

for the quasipotential and

$$\Psi_{\sigma_1\sigma_2}(r,\vec{n}) = \frac{1}{(2\pi)^3}\int d\Omega_p\, \xi(\vec{p};\vec{r})\Psi_{\sigma_1\sigma_2}(\vec{p}), \tag{3.2}$$

for the WF. The integration measure is

$$d\Omega_p \equiv d^3\vec{p}/\sqrt{1+\vec{p}^2/M^2} \tag{3.3}$$

It is the invariant measure on the hyperboloid, $p_0^2 - \vec{p}^2 = M^2$. The system of functions $\xi(\vec{p};\vec{n},r)$ is the complete orthogonal system of functions in the Lobachevsky space,

$$\xi(\vec{p};\vec{n};r) = (\frac{p_0 - \vec{p}\vec{n}}{M})^{-1-irM}. \tag{3.4}$$

The physical meaning of the parameter r is discussed in details in ref. [31].

As a result of carring out this transformation to the RCR we arrive at the following quasipotential:

$$V(r,\vec{n};p_0,\vec{p}) = V_1(r,p_0) + V_2(r,\vec{n};p_0,\vec{p}), \tag{3.5}$$

with

$$\begin{aligned} V_1(r;p_0) = -3f^2 \Bigg\{ & (8p_0^2 - 4M^2)V_{Yuk}(r) + (\frac{3p_0^2}{M^2} - 1)\frac{1}{r}\delta(r^2 + \frac{1}{M^2}) + \\ + \quad & \frac{p_0^2}{M^2}\frac{1}{r}\delta(r^2+\frac{4}{M^2}) + \frac{2\vec{p}^2}{M^2}\left[B(r) + \frac{1}{3}\frac{1}{r}\delta(r^2+\frac{1}{M^2})\right] - \\ - \quad & (\vec{S}_1\vec{S}_2)\left[\frac{2\vec{p}^2}{M^2}B(r) + \frac{1}{6}\frac{1}{r}\delta(r^2+\frac{1}{M^2}) + \frac{1}{2}\frac{1}{r}\delta(r^2+\frac{4}{M^2})\right]\Bigg\}, \end{aligned} \tag{3.6}$$

and

$$\begin{aligned} V_2(r,\vec{n};p_0,\vec{p}) = -3f^2 \Bigg\{ & \frac{2}{M^2}(\vec{S}_1\vec{p})(\vec{S}_2\vec{p})B(r) - S_{12}B(r) + \\ + \quad & \frac{3}{M^2}\left[(\vec{S}_1\vec{L})(\vec{S}_2\vec{L}) + (\vec{S}_2\vec{L})(\vec{S}_1\vec{L})\right]\frac{1}{r^2}B(r) - \frac{6}{M^2}(\vec{p}\vec{n})^2B(r) + \\ + \quad & \frac{2ip_0}{M}(\vec{p}\vec{n})\left[4rA(r) - \frac{1}{M^2}C(r)\right] - (\vec{S}\vec{L})\left[\frac{4p_0}{M}A(r) + \right. \\ + \quad & \left.\frac{6i}{M^2}(\vec{p}\vec{n})\frac{1}{r}B(r) - \frac{p_0}{M^3}\frac{1}{r}C(r)\right]\Bigg\}, \end{aligned} \tag{3.7}$$

where $\vec{S} = \vec{S}_1 + \vec{S}_2$, $\vec{L} = [\vec{p} \times \vec{r}]$, $S_{12} = 3(\vec{S}_1\vec{n})(\vec{S}_2\vec{n}) - (\vec{S}_1\vec{S}_2)$.

Next,

$$V_{Yuk}(r) = \frac{1}{4\pi r} cth(rM\pi), \tag{3.8}$$

and

$$\begin{aligned} A(r) &= \frac{1}{r(r+(i/M))} V_{Yuk}(r), & (3.9)\\ B(r) &= \frac{1}{(r+(i/M))(r+(2i/M))} V_{Yuk}(r), & (3.10)\\ C(r) &= \frac{1}{(i/M)} \frac{(r-(i/M))}{r+(i/M)} \frac{1}{r} \delta(r^2 + \frac{4}{M^2}) - \frac{M^2}{4r}\delta(r). & (3.11) \end{aligned}$$

From the above equations we can see that the quasipotential is separated into two parts, $V_1(r;p_0)$, which does not depend on the direction of the "relativistic coordinate" vector, and $V_2(r,\vec{n};p_0,\vec{p})$, which depends on $\vec{n}$ over the structure $(\vec{p}\vec{n})$, $[\vec{p}\times\vec{n}]$, S_{12}.

After the transformation of the quasipotential equations:

$$(\mathcal{M} - 2p_0)\Psi_{\sigma_1\sigma_2}(\vec{p}) = (2\pi)^{-3} \sum_{\nu_1\nu_2} \int d\Omega_k V^{\nu_1\nu_2}_{\sigma_1\sigma_2}(\vec{k},\vec{p})\Psi_{\nu_1\nu_2}(\vec{k}), \tag{3.12}$$

in the RCR by means of the formulae (3.1,3.2), we can be convinced that $V_1(r;p_0)$ describes the local type of interaction and $V_2(r,\vec{n};p_0,\vec{p})$ enters into equation (3.13) by the non-local way.

$$\begin{aligned} (\mathcal{M} - 2\hat{H})\Psi_{\sigma_1\sigma_2}(\vec{r}) &= \sum_{\nu_1\nu_2} V^{\nu_1\nu_2}_{1_{\sigma_1\sigma_2}}(r;p_0)\Psi_{\nu_1\nu_2}(\vec{r}) + \int d^3\vec{r}_1 \sum_{\nu_1\nu_2} \int d\Omega_p \times \\ &\times \quad \xi(\vec{p};\vec{n},r)\xi(\vec{p};\vec{n}_1,r_1) V^{\nu_1\nu_2}_{2_{\sigma_1\sigma_2}}(r_1,\vec{n}_{1\Lambda p};p_0,\vec{p})\Psi_{\nu_1\nu_2}(\vec{r}_1), & (3.13) \end{aligned}$$

where the unit vector is [32]

$$\vec{n}_{\Lambda p} = [M\vec{n} - \vec{p}(1 - \frac{\vec{p}\vec{n}}{p_0 + M})]/(p - \vec{p}\vec{n}). \tag{3.14}$$

It is still possible to localize the spin-orbit part, some terms of the tensor interaction and some other terms entering in $V_2(r_1, \vec{n}_{1\Lambda_p}; p_0, \vec{p})$.

We have the following structure in Eq. (3.13):

$$\begin{aligned} (\vec{p}\cdot\vec{n}_{1\Lambda p}) &= M^2/(p_0 - \vec{p}\vec{n}_1) - p_0, & (3.15)\\ [\vec{p}\times\vec{n}_{1\Lambda p}] &= M\,[\vec{p}\times\vec{n}_1]\,/(p_0 - \vec{p}\vec{n}_1), & (3.16) \end{aligned}$$

and

$$(\vec{S}_1\vec{n}_{1\Lambda p})(\vec{S}_2\vec{n}_{1\Lambda p}) = Z_1^T + Z_2^T, \tag{3.17}$$

where

$$\begin{aligned} Z_1^T &= M^2(\vec{S}_1\vec{n}_1)(\vec{S}_2\vec{n}_1)/(p_0 - \vec{p}\vec{n}_1), & (3.18)\\ Z_2^T &= \frac{M^2}{(p_0 - \vec{p}\vec{n}_1)^2}\{-\frac{1}{M}[(\vec{S}_1\vec{n}_1)(\vec{S}_2\vec{p}) + (\vec{S}_1\vec{p})(\vec{S}_2\vec{n}_1)] \times \\ &\times \quad (1 - \frac{\vec{p}\vec{n}_1}{p_0 + M}) + \frac{1}{M^2}(\vec{S}_1\vec{p})(\vec{S}_2\vec{p})(1 - \frac{\vec{p}\vec{n}_1}{p_0 + M})^2\}. & (3.19) \end{aligned}$$

The localization procedure for (3.15,3.16) and the first term of Z_1^T is produced by means of the following equation:

$$\exp(\frac{i}{M}\frac{\partial}{\partial r_1})\xi^*(\vec{p};\vec{n}_1,r_1) = \frac{M}{p_0-\vec{p}\vec{n}_1}\xi^*(\vec{p};\vec{n}_1,r_1). \tag{3.20}$$

As a result we obtain

$$(\mathcal{M}-2\hat{H})\Psi_{\sigma_1\sigma_2}(\vec{r}) = \sum_{\nu_1\nu_2}\hat{V}^{\nu_1\nu_2}_{\sigma_1\sigma_2}(\vec{r};p_0,\vec{p})\Psi_{\nu_1\nu_2}(\vec{r}). \tag{3.21}$$

The quasipotential of Eq. (3.13) is presented as a sum of six component:

$$\hat{V}(\vec{r};p_0,\vec{p}) = \hat{V}_C + \hat{V}_{LS}(\vec{L}\vec{S}) + \hat{V}_S(\vec{S}_1\vec{S}_2) + \hat{V}_T S_{12} + \hat{V}_{LL}L_{12} + \hat{V}_{Sp}P_{12}, \tag{3.22}$$

where

$$\begin{aligned} L_{12} &= \frac{1}{2}\{(\vec{S}_1\vec{L})(\vec{S}_2\vec{L}) + (\vec{S}_2\vec{L})(\vec{S}_1\vec{L})\}, && (3.23)\\ P_{12} &= (\vec{S}_1\vec{p})(\vec{S}_2\vec{p}); && (3.24)\end{aligned}$$

and

$$\begin{aligned} \hat{V}_C &= -3f^2\left\{(8p_0-4M^2)V_{Yuk}(r) + (\frac{3p_0^2}{M^2}-1)\frac{1}{r}\delta(r^2+\frac{1}{M^2}) + \frac{p_0^2}{M^2}\frac{1}{r}\delta(r^2+\frac{4}{M^2}) - \right.\\ &+ 2\frac{\vec{p}^2}{M^2}\left[B(r) + \frac{1}{3}\frac{1}{r}\delta(r^2+\frac{1}{M^2})\right] - 6\left[\frac{(r-\frac{2i}{M})^2}{r^2}\exp(-\frac{2i}{M}\frac{\partial}{\partial r}) - \frac{2p_0}{M}\frac{(r-\frac{i}{M})^2}{r^2}\times\right.\\ &\times \left.\exp(-\frac{i}{M}\frac{\partial}{\partial r}) + \frac{p_0^2}{M^2}\right]B(r) + \frac{2ip_0}{M}\left[M\frac{(r-\frac{i}{M})^2}{r^2}\exp(-\frac{i}{M}\frac{\partial}{\partial r}) - p_0\right]\times\\ &\times \left.\left[4rA(r) - \frac{1}{M^3}C(r)\right]\right\}, && (3.25)\\ \hat{V}_{LS} &= -3f^2\left\{\frac{p_0}{M}\frac{r^2}{(r+(i/M))^2}\left[4A(r) - \frac{1}{M^2}\frac{1}{r}C(r)\right] + \right.\\ &+ \left.\frac{6i}{M^2}\left[M\frac{r(r-(i/M))}{(r+(i/M))^2}\exp(-\frac{i}{M}\frac{\partial}{\partial r}) - p_0\right]\frac{r}{(r+(i/M))^2}B(r)\right\}, && (3.26)\\ \hat{V}_S &= -3f^2\left\{-\left[\frac{(r-(2i/M))^2}{r^2}\exp(-\frac{2i}{M}\frac{\partial}{\partial r}) + \frac{2\vec{p}^2}{M^2} - 1\right]B(r) - \right.\\ &- \left.\frac{1}{6}\frac{1}{r}\delta(r^2+\frac{1}{M^2}) - \frac{1}{2}\frac{1}{r}\delta(r^2+\frac{4}{M^2})\right\}, && (3.27)\\ \hat{V}_T &= -3f^2\left\{-\frac{(r-(2i/M))^2}{r^2}\exp(-\frac{2i}{M}\frac{\partial}{\partial r})B(r)\right\}, && (3.28)\\ \hat{V}_{LL} &= -3f^2\left\{\frac{3}{M^2}\frac{r}{(r+(i/M))(r+(2i/M))^2}B(r)\right\}, && (3.29)\\ \hat{V}_{Sp} &= -3f^2\{\frac{2}{M^2}B(r)\}. && (3.30)\end{aligned}$$

It is noted that we have assumed the corresponding finite-difference operators in p_0 and $\vec{p}$ for these calculations. However, the use of eigenvalues instead of these operators is a good approximation, which does not reduce to the non-relativistic models.

After carrying out the partial-wave expansion for the WF

$$\Psi^{(S)}(\vec{r};\sigma) = \frac{4\pi}{r}\sum_{J\ell M} R_{J\ell S}(r)\left\{\Omega^{*(S)}_{J\ell M}(\vec{n})\right\}_{\sigma}, \tag{3.31}$$

the three-dimensional quasipotential equation is rewritten in the system of radial equations for $S=0$, $S=1$ and $S=2$

$$(\mathcal{M}-2\hat{H}_{\ell})R_{J\ell S}(r) = \sum_{\ell' S'} \hat{V}^{J}_{\ell S,\ell' S'}(r;p_0,\vec{p})R_{J\ell' S'}(r). \tag{3.32}$$

Here,

$$\hat{H}_{\ell} = M\, cosh(\frac{i}{M}\frac{\partial}{\partial r}) + \frac{i}{r} sinh(\frac{i}{M}\frac{\partial}{\partial r}) + \frac{\ell(\ell+1)}{2Mr^2}\exp(\frac{i}{M}\frac{\partial}{\partial r}), \tag{3.33}$$

and

$$\hat{V}^{J}_{\ell' S',\ell S}(r;p_0,\vec{p}) = \int d\omega_{\vec{n}}\Omega^{*(S')}_{J\ell' M}(\vec{n})\hat{V}(\vec{r};p_0,\vec{p})\Omega^{(S)}_{J\ell M}(\vec{n}). \tag{3.34}$$

One can check that due to the complete accordance of the relativistic spin structure of the quasipotential (3.22) to the spin structures which are used in non-relativistic models, the matrix elements (3.34) have the same form as when non-relativistic interaction of two vector particles is considered. The only differences are in the internal expressions of $\hat{V}_C$, $\hat{V}_T$, $\hat{V}_{SL}$, $\hat{V}_S$, $\hat{V}_{LL}$ and $\hat{V}_{Sp}$.

Thus, we have for a singlet state

$$(\mathcal{M}-2\hat{H}_{\ell=J})R_{\ell=J}(r) = \left(\hat{V}_C - 2\hat{V}_S - \frac{2}{3}J(J+1)\hat{V}_{LL}\right)R_{\ell=J}(r); \tag{3.35}$$

for a triplet state

$$\begin{aligned}(\mathcal{M}-2\hat{H}_{\ell=J+1})R_{\ell=J+1}(r) &= \left(\hat{V}_C - (J+2)\hat{V}_{SL} - \hat{V}_S - \frac{J+2}{2J+1}\hat{V}_T + \right.\\ &\left. + \frac{J+2}{2}\hat{V}_{LL}\right)R_{\ell=J+1}(r) + \frac{3\sqrt{J(J+1)}}{2J+1}\hat{V}_T R_{\ell=J-1}(r),\end{aligned} \tag{3.36}$$

$$(\mathcal{M}-2\hat{H}_{\ell=J})\quad R_{\ell=J}(r) = \left(\hat{V}_C - \hat{V}_{LS} - \hat{V}_S + \hat{V}_T + [\frac{1}{2} - J(J+1)]\hat{V}_{LL}\right)R_{\ell=J}(r), \tag{3.37}$$

$$\begin{aligned}(\mathcal{M}-2\hat{H}_{\ell=J-1})R_{\ell=J-1}(r) &= \left(\hat{V}_C + (J-1)\hat{V}_{LS} - \hat{V}_S - \frac{J-1}{2J+1}\hat{V}_T - \right.\\ &\left. - \frac{J-1}{2}\hat{V}_{LL}\right)R_{\ell=J-1}(r) + \frac{3\sqrt{J(J+1)}}{2J+1}\hat{V}_T R_{\ell=J+1}(r);\end{aligned} \tag{3.38}$$

for 5-plet state

$$(\mathcal{M}-2\hat{H}_{\ell=J+2})R_{\ell=J+2}(r) = \left(\hat{V}_C - 2(J+3)\hat{V}_{LS} + \hat{V}_S - \frac{2(J+3)}{2J+3}\hat{V}_T + \right.$$

$$+(J+3)^2\hat{V}_{LL}\Big)\, R_{\ell=J+2}(r) + \frac{\sqrt{6J(J+2)}}{2J+1}\frac{\sqrt{(2J-1)(2J+1)}}{2J+3}\hat{V}_T R_{\ell=J}(r), \tag{3.39}$$

$$(\mathcal{M}-2\hat{H}_{\ell=J+1})R_{\ell=J+1}(r) = \Big(\hat{V}_C - (J+4)\hat{V}_{LS} + \hat{V}_S + \frac{J-4}{2J+1}\hat{V}_T +$$

$$+\frac{3J+8}{2}\hat{V}_{LL}\Big)\, R_{\ell=J+1}(r) + \frac{3\sqrt{(J-1)(J+2)}}{2J+1}\hat{V}_T R_{\ell=J-1}(r), \tag{3.40}$$

$$(\mathcal{M}-2\hat{H}_{\ell=J})R_{\ell=J}(r) = \frac{\sqrt{6J(J+2)}}{2J+1}\frac{\sqrt{(2J-1)(2J+1)}}{2J+3}\hat{V}_T R_{\ell=J+2}(r) +$$

$$+\left(\hat{V}_C - 3\hat{V}_{LS} + \hat{V}_S + \frac{(2J-3)(2J+5)}{(2J-1)(2J+3)}\hat{V}_T + \left[\frac{5}{2} - \frac{1}{3}J(J+1)\right]\hat{V}_{LL}\right)\times$$

$$\times R_{\ell=J}(r) + \frac{\sqrt{6(J-1)(J+1)}}{2J+1}\frac{\sqrt{(2J+1)(2J+3)}}{2J-1}\hat{V}_T R_{\ell=J-2}(r), \tag{3.41}$$

$$(\mathcal{M}-2\hat{H}_{\ell=J-1})R_{\ell=J-1}(r) = \Big(\hat{V}_C + (J-3)\hat{V}_{LS} + \hat{V}_S + \frac{J+5}{2J+1}\hat{V}_T -$$

$$-\frac{3J-5}{2}\hat{V}_{LL}\Big)\, R_{\ell=J-1}(r) + \frac{3\sqrt{(J-1)(J+2)}}{2J+1}\hat{V}_T R_{\ell=J+1}(r), \tag{3.42}$$

$$(\mathcal{M}-2\hat{H}_{\ell=J-2})R_{\ell=J-2}(r) = \Big(\hat{V}_C + 2(J-2)\hat{V}_{LS} + \hat{V}_S - \frac{2(J-2)}{2J-1}\hat{V}_T +$$

$$+(J-2)^2\hat{V}_{LL}\Big)\, R_{\ell=J-2}(r) + \frac{\sqrt{6(J-1)(J+1)}}{2J+1}\frac{\sqrt{(2J+1)(2J+3)}}{2J-1}\hat{V}_T R_{\ell=J}(r). \tag{3.43}$$

The two-gluon bound states are to have the positive C- parity. Therefore, we have restrictions on J in (3.35-3.43).

Finally, let us remark that we have neglected the last term of (3.22) following the authors of ref. [33].

4 Quasiclassical condition for a quantization in RCR

The non-tied partial-wave equations can be re-written in the following form:

$$[cosh(i\bar{\lambda}\frac{\partial}{\partial r}) + \frac{i\bar{\lambda}}{r}sinh(i\bar{\lambda}\frac{\partial}{\partial r}) + \frac{\bar{\lambda}^2\ell(\ell+1)}{2r^2}\exp(i\bar{\lambda}\frac{\partial}{\partial r}) - X(r)]R_\ell(r) = 0, \tag{4.1}$$

$$X(r) = \frac{W - V(r)}{2M}, \tag{4.2}$$

where $W = 2M + E_b$ is the mass of bound state, M is the gluon mass, E_b is the binding energy. We consider $V(r)$ as a sum of the potential V_{conf} describing the confinement of a gluon in a meson and the corresponding matrix elements of the quasipotential in the system of the Eqs. (3.32). In the expressions (3.25-3.30) it is possible to neglect the

small image additions which are proportional to $i/M(= i\hbar/Mc)$, because the bound state spectrum usually forms on distances of the order $r >> \bar{\lambda} = \hbar/Mc$.

The quasiclassical condition for a quantization of the two-particle relativistic systems has the following form [19]:

$$\int_{r_-}^{r_+} dr' Arch X_\Lambda(r') = \bar{\lambda}\pi(n + \frac{1}{2}), \tag{4.3}$$

where

$$X_\Lambda(r) = X(r)[1 + (\Lambda\frac{\bar{\lambda}}{r})^2]^{-1/2}, \quad \Lambda = \ell + 1, \tag{4.4}$$

and the integration limits are determined from the equation

$$X_\Lambda(r_\pm) = 1. \tag{4.5}$$

If restrict ourselves to the cases of the simplest potentials, namely, $V(r) = \sigma r$ and $V(r) = \sigma r^2$, the quantization condition is shown in ref. [34] to take the form of (4.6) and (4.7), correspondingly,

$$\chi\, cosh\, \chi - sinh\, \chi = \frac{\sigma}{2Mc^2}\bar{\lambda}\pi(n + \frac{\ell}{2} + \frac{3}{4}), \tag{4.6}$$

$$2\sqrt{cosh\, \chi + 1}\,[K(tanh\, \chi/2) - E(tanh\, \chi/2)] = \sqrt{\frac{\sigma}{2Mc^2}}\bar{\lambda}\pi(n + \frac{\ell}{2} + \frac{3}{4}), \tag{4.7}$$

where $cosh\, \chi = W/2M$. Following this technique, the energy levels of gluonium states can be obtained. The numerical results for the case of the confinement potential are presented in ref. [26]. Investigation of the influence of the Coulomb term as well as the relativistic corrections are in progress.

5 Conclusions

In the present work the formalism for consideration of the two vector particles, gluons, has been constructed on the basis of the $2(2S + 1)$- dimensional description of the WF. The form of the relativistic two-particle single-time quasipotential equation was found. It turned out that the quasipotential of this equation coincided with the quasipotential for interaction of two spinor particles in the second order of perturbation theory. This fact shows the advantages of the $2(2S + 1)$- component formalism.

It is possible to employ the relativistic generalization of the WKB method to find the gluonium spectrum in the case when the total spin of the system is equal to zero. In following works we will use the obtained system of the partial-wave equations for investigation of the contributions of the spin-spin and spin-orbit interactions to the energy of the gluonium states.

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